Gribov horizon in the presence of dynamical mass generation in Euclidean Yang-Mills theories in the Landau gauge

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Abstract

The infrared behavior of the gluon and ghost propagators is analyzed in Yang-Mills theories in the presence of dynamical mass generation in the Landau gauge. By restricting the domain of integration in the path-integral to the Gribov region Ω, the gauge propagator is found to be suppressed in the infrared, while the ghost propagator is enhanced.

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1 Introduction.

The possibility that gluons might acquire a mass through a dynamical mechanism is receiving renewed interest in the last few years. Although a fully gauge invariant framework for the dynamical mass generation in Yang-Mills theories is not yet available, the number of gauges displaying this interesting phenomenon is getting considerably large.

A dynamical gluon mass has been introduced in the light-cone gauge [1] in order to obtain estimates for the spectrum of the glueballs. It has been discussed in the Coulomb gauge in [2], where the presence of a nonvanishing condensate $\langle A_i^a A_i^a \rangle$ in the operator product expansion for the two-point gauge correlation function has been pointed out. More recently, the condensate $\left\langle A_\mu^a A_\mu^a \right\rangle$ has been investigated in the Landau gauge in [3, 4], where it has been proven to account for the discrepancy observed in the two- and three-point correlation functions between the perturbative theory and the lattice results. A renormalizable effective potential for the condensate $\left\langle A_\mu^a A_\mu^a \right\rangle$ in pure Yang-Mills theory in the Landau gauge has been constructed and evaluated in analytic form up to two-loop order in [5, 6]. This result shows that the vacuum of Yang-Mills theory favors the formation of a nonvanishing condensate $\left\langle A_\mu^a A_\mu^a \right\rangle$, which lowers the vacuum energy and provides a dynamical gluon mass, which turns out to be of the order of $\approx 500MeV$. The inclusion of massless quarks has been worked out in [7]. We remind here that lattice simulations of the gluon propagator in the Landau gauge have reported a gluon mass $m \approx 600MeV$ [8]. Concerning other gauges, the occurrence of the condensate $\left\langle A_\mu^a A_\mu^a \right\rangle$ and of the related dynamical gluon mass has been established in the linear covariant gauges in [9, 10]. These results can be generalized to a class of nonlinear covariant gauges. Here, the mixed gluon-ghost condensate $\left\langle \frac{1}{2} A_\mu^a A_\mu^a + \xi \overline{c}^a c^a \right\rangle$ has to be considered [11], with ξ the gauge parameter. A renormalizable effective potential for this condensate has been obtained in the Curci-Ferrari [12] and Maximal Abelian gauges [13], resulting in a dynamical mass generation. In the latter case, lattice simulations [14] had already given evidences of a nonvanishing mass for the off-diagonal gluons. Moreover, a gluon mass has been reported in lattice simulations in the Laplacian gauge [15]. Also, it is part of the so-called Kugo-Ojima criterion for color confinement [16] and, as discussed in [17], it proves to be useful in order to account for the data obtained on the radiative decays of heavy quarkonia systems.

In this work we pursue the study of the dynamical mass generation in Euclidean Yang-Mills theory in the Landau gauge. We attempt at incorporating the nonperturbative effects related to the Gribov horizon [18], the aim being that of investigating the infrared behavior of the gluon and ghost propagators in presence of the dynamical mass generation. These propagators have been studied to a great extent by several groups through lattice simulations [8, 19, 20, 21, 22] in the Landau gauge, which have confirmed that the gluon propagator is suppressed in the infrared region while the ghost propagator is enhanced, being in fact more singular than the perturbative behavior $\approx 1/k^2$. Such behavior of the gluon and ghost propagators was already found by Gribov in [18], where it arises as a consequence of the restriction of the domain of integration in the path-integral to the region Ω whose boundary $\partial\Omega$ is the first Gribov horizon, where the first vanishing eigenvalue of the Faddeev-Popov operator, $-\partial_\mu \left(\partial_\mu \delta^{ab} + g f^{acb} A_\mu^c \right)$, appears. This restriction is necessary due to the existence of the Gribov copies, which imply that the Landau condition, $\partial_\mu A_\mu^a = 0$, does not uniquely fix the gauge. The infrared suppression of the gluon propagator and the enhancement of the ghost propagator have also been derived in [23], where the restriction to the region Ω has been implemented by a Boltzmann factor through

the introduction of a horizon function. Recently, the authors of [24, 25, 26, 27] have analyzed the behavior of the gluon and ghost propagators in the Landau gauge within the Schwinger-Dyson framework, also obtaining that the gluon propagator is suppressed while the ghost propagator is enhanced.

Concerning now the gluon and ghost propagators in the presence of a dynamical mass generation, we shall proceed by following Gribov's original suggestion, which amounts to implement the restriction to Ω as a no-pole condition for the two-point ghost function [18]. We shall be able to show that the gluon and ghost propagators are suppressed and enhanced, respectively, and this in the presence of a dynamical gluon mass. This behavior is in agreement with that found in [18, 23, 24, 25, 26, 27].

This work is organized as follows. In Sect.2 we briefly review the properties of the Lagrangian accounting for the dynamical gluon mass generation in the Landau gauge. In Sect.3 we implement the restriction of the domain of integration in the path-integral to the region Ω. The ensuing modifications of the gauge propagator due to both the Gribov horizon and dynamical gluon mass are worked out. Sect.4 is devoted to the analysis of the infrared behavior of the ghost propagator. Some further remarks are collected in the Conclusion.

2 Dynamical mass generation in the Landau gauge.

The dynamical mass generation in the Landau gauge is described by the following action [5]

$$S(A,\sigma) = S_{YM} + S_{GF+FP} + S_\sigma \ , \tag{2.1}$$

where S_{YM}, S_{GF+FP} are the Yang-Mills and the gauge fixing terms

$$S_{YM} = \frac{1}{4}\int d^4x F^a_{\mu\nu}F^a_{\mu\nu} \ , \tag{2.2}$$

$$S_{GF+FP} = \int d^4x \left(b^a\partial_\mu A^a_\mu + \overline{c}^a \partial_\mu D^{ab}_\mu c^b \right) \ , \tag{2.3}$$

with b^a being the Lagrange multiplier enforcing the Landau gauge condition, $\partial_\mu A^a_\mu = 0$, and $\overline{c}^a$, c^a denoting the Faddeev-Popov ghosts. The color index a refers to the adjoint representation of the gauge group $SU(N)$. The term S_σ in eq.(2.1) contains the auxiliary scalar field σ and reads

$$S_\sigma = \int d^4x \left(\frac{\sigma^2}{2g^2\zeta} + \frac{1}{2}\frac{\sigma}{g\zeta}A^a_\mu A^a_\mu + \frac{1}{8\zeta}\left(A^a_\mu A^a_\mu\right)^2 \right) . \tag{2.4}$$

The introduction of the auxiliary field σ allows to study the condensation of the local operator $A^a_\mu A^a_\mu$. In fact, as shown in [5], the following relation holds

$$\langle\sigma\rangle = -\frac{g}{2}\left\langle A^a_\mu A^a_\mu \right\rangle \ . \tag{2.5}$$

The dimensionless parameter ζ in expression (2.4) is needed to account for the ultraviolet divergences present in the vacuum correlation function $\langle A^2(x)A^2(y)\rangle$. For the details of the renormalizability properties of the local operator $A^a_\mu A^a_\mu$ in the Landau gauge we refer to [28, 29]. Expression (2.1) is left invariant by the following BRST transformations

$$\begin{aligned}
sA^a_\mu &= -D^{ab}_\mu c^b = -\left(\partial_\mu c^a + gf^{abc}A^b_\mu c^c\right) \ , \\
sc^a &= \frac{1}{2}gf^{abc}c^b c^c \ ,
\end{aligned}$$

$$
\begin{aligned}
s\overline{c}^a &= b^a \ , \\
sb^a &= 0 \ , \\
s\sigma &= gA_\mu^a \partial_\mu c^a \ ,
\end{aligned}
\tag{2.6}
$$

and

$$
sS(A, \sigma) = 0 \ . \tag{2.7}
$$

Notice that, from the relation

$$
A_\mu^a \partial_\mu c^a = -\frac{1}{2} s \left(A_\mu^a A_\mu^a \right) \ , \tag{2.8}
$$

it follows that the BRST operator is nilpotent. The action $S(A, \sigma)$ is the starting point for constructing a renormalizable effective potential $V(\sigma)$ for the auxiliary field σ, obeying the renormalization group equations. The output of the higher loop computations done in [5, 7] shows that the minimum of $V(\sigma)$ occurs for a nonvanishing vacuum expectation value of the auxiliary field, *i.e.* $\langle \sigma \rangle \neq 0$. In particular, from expression (2.1), the first order induced dynamical gluon mass is found to be

$$
m^2 = \frac{g \langle \sigma \rangle}{\zeta_0} \ , \tag{2.9}
$$

where ζ_0 is the first contribution to the parameter ζ [5], given by

$$
\begin{aligned}
\zeta &= \frac{\zeta_0}{g^2} + \zeta_1 + O(g^2) \ , \\
\zeta_0 &= \frac{9}{13} \frac{(N^2 - 1)}{N} \ .
\end{aligned}
\tag{2.10}
$$

We remind here that, in the Landau gauge, the Faddeev-Popov ghosts $\overline{c}^a$, c^a remain massless, due to the absence of mixing between the composite operators $A_\mu^a A_\mu^a$ and $\overline{c}^a c^a$. This stems from additional Ward identities present in the Landau [29] and in the covariant linear gauges [9], which forbid the appearance of the term $\overline{c}^a c^a$.

3 Infrared behavior of the gluon propagator.

3.1 Restriction to the region Ω.

In the previous section we have reviewed the properties of the action $S(A, \sigma)$ which accounts for the dynamical mass generation. However, it is worth underlining that the action $S(A, \sigma)$ leads to a partition function

$$
\mathcal{Z} = \mathcal{N} \int DAD\sigma \, \delta(\partial A^a) \det \left(-\partial_\mu \left(\partial_\mu \delta^{ab} + g f^{acb} A_\mu^c \right) \right) e^{-(S_{YM} + S_\sigma)} \ , \tag{3.11}
$$

which is still plagued by the Gribov copies, which affect the Landau gauge. It might be useful to notice here that the action $(S_{YM} + S_\sigma)$ is left invariant by the local gauge transformations

$$
\begin{aligned}
\delta A_\mu^a &= -D_\mu^{ab} \omega^b \ , \\
\delta \sigma &= g A_\mu^a \partial_\mu \omega^a \ ,
\end{aligned}
\tag{3.12}
$$

$$
\delta \left(S_{YM} + S_\sigma \right) = 0 \ . \tag{3.13}
$$

As a consequence of the existence of Gribov copies, the domain of integration in the path-integral should be restricted further. We shall follow here Gribov's proposal to restrict the domain of integration to the region Ω [18]. Expression (3.11) is thus replaced by

$$\mathcal{Z} = \mathcal{N} \int DAD\sigma\ \delta(\partial A^a) \det\left(-\partial_\mu \left(\partial_\mu \delta^{ab} + g f^{acb} A^c_\mu\right)\right) e^{-(S_{YM}+S_\sigma)}\ \mathcal{V}(\Omega)\ , \tag{3.14}$$

where $\mathcal{V}(\Omega)$ implements the restriction to Ω. The factor $\mathcal{V}(\Omega)$ can be accommodated for by requiring that the two-point connected ghost function $\mathcal{G}(k,A)$ has no poles for a given nonvanishing value of the momentum k [18]. This condition can be understood by recalling that the region Ω is defined as the set of all transverse gauge connections $\left\{A^a_\mu\right\}$, $\partial_\mu A^a_\mu = 0$, for which the Faddeev-Popov operator is positive definite, *i.e.* $-\partial_\mu \left(\partial_\mu \delta^{ab} + g f^{acb} A^c_\mu\right) > 0$. This implies that the inverse of the Faddeev-Popov operator $\left[-\partial_\mu \left(\partial_\mu \delta^{ab} + g f^{acb} A^c_\mu\right)\right]^{-1}$, and thus $\mathcal{G}(k,A)$, can become large only when approaching the horizon $\partial\Omega$, which corresponds in fact to $k=0$ [18]. The quantity $\mathcal{G}(k,A)$ can be evaluated order by order in perturbation theory. Repeating the same calculation of [18], we find that, up to the second order

$$\mathcal{G}(k,A) \approx \frac{1}{k^2}\frac{1}{1-\rho(k,A)}\ , \tag{3.15}$$

with $\rho(k,A)$ given by

$$\rho(k,A) = \frac{g^2}{3}\frac{N}{N^2-1}\frac{1}{V}\frac{k_\mu k_\nu}{k^2}\sum_q \frac{1}{(k-q)^2}\left(A^a_\lambda(q)A^a_\lambda(-q)\right)\left(\delta_{\mu\nu} - \frac{q_\mu q_\nu}{q^2}\right)\ , \tag{3.16}$$

and V being the space-time volume. According to [18], the no-pole condition for $\mathcal{G}(k,A)$ reads

$$\begin{aligned} \rho(0,A) &< 1\ , \\ \rho(0,A) &= \frac{g^2}{4}\frac{N}{N^2-1}\frac{1}{V}\sum_q \frac{1}{q^2}\left(A^a_\lambda(q)A^a_\lambda(-q)\right)\ . \end{aligned} \tag{3.17}$$

Therefore, for the factor $\mathcal{V}(\Omega)$ in eq.(3.14) we have

$$\mathcal{V}(\Omega) = \theta(1-\rho(0,A))\ , \tag{3.18}$$

where $\theta(x)$ stands for the step function*.

3.2 The gauge propagator.

In order to discuss the gauge propagator, it is sufficient to retain only the quadratic terms in expression (3.14) which contribute to the two-point correlation function $\left\langle A^a_\mu(k)A^b_\nu(-k)\right\rangle$. Expanding around the nonvanishing vacuum expectation value of the auxiliary field, $\langle\sigma\rangle \neq 0$, and making use of the integral representation for the step function

$$\theta(1-\rho(0,A)) = \int_{-i\infty+\varepsilon}^{i\infty+\varepsilon} \frac{d\eta}{2\pi i\eta} e^{\eta(1-\rho(0,A))}\ , \tag{3.19}$$

we get

$$\begin{aligned} \mathcal{Z}_{\text{quadr}} &= \mathcal{N}\int DA \frac{d\eta}{2\pi i\eta} e^{\eta(1-\rho(0,A))}\delta(\partial A^a) e^{-\left(\frac{1}{4}\int d^4x((\partial_\mu A^a_\nu - \partial_\mu A^a_\nu)^2 + \frac{1}{2}m^2\int d^4x\left(A^a_\mu A^a_\mu\right)\right)} \\ &= \mathcal{N}\int DA \frac{d\eta}{2\pi i\eta} e^{\eta} e^{-\frac{1}{2}\sum_q A^a_\mu(q)\mathcal{Q}^{ab}_{\mu\nu}A^b_\nu(-q)}\ , \end{aligned} \tag{3.20}$$

*$\theta(x) = 1$ if $x > 0$, $\theta(x) = 0$ if $x < 0$.

with

$$\mathcal{Q}^{ab}_{\mu\nu} = \left(\left(q^2 + m^2\right)\delta_{\mu\nu} + \left(\frac{1}{\alpha} - 1\right) q_\mu q_\nu + \frac{\eta N g^2}{N^2-1}\frac{1}{2V}\frac{1}{q^2}\delta_{\mu\nu} \right)\delta^{ab} \ , \tag{3.21}$$

where the limit $\alpha \to 0$ has to be taken at the end in order to recover the Landau gauge. Integrating over the gauge field, one has

$$\mathcal{Z}_{\rm quadr} = \mathcal{N} \int \frac{d\eta}{2\pi i \eta} e^{\eta} \left(\det \mathcal{Q}^{ab}_{\mu\nu}\right)^{-\frac{1}{2}} = \mathcal{N} \int \frac{d\eta}{2\pi i} e^{f(\eta)} \ , \tag{3.22}$$

where $f(\eta)$ is given by

$$f(\eta) = \eta - \log \eta - \frac{3}{2}(N^2-1)\sum_q \log\left(q^2 + m^2 + \frac{\eta N g^2}{N^2-1}\frac{1}{2V}\frac{1}{q^2}\right) \ . \tag{3.23}$$

Following [18], the expression (3.22) can be now evaluated at the saddle point, namely

$$\mathcal{Z}_{\rm quadr} \approx e^{f(\eta_0)} \ , \tag{3.24}$$

where η_0 is determined by the minimum condition

$$1 - \frac{1}{\eta_0} - \frac{3}{4}\frac{Ng^2}{V}\sum_q \frac{1}{\left(q^4 + m^2 q^2 + \frac{\eta_0 N g^2}{N^2-1}\frac{1}{2V}\right)} = 0 \ . \tag{3.25}$$

Taking the thermodynamic limit, $V \to \infty$, and setting [18]

$$\gamma^4 = \frac{\eta_0 N g^2}{N^2-1}\frac{1}{2V} \quad , \quad V \to \infty \ , \tag{3.26}$$

we get the gap equation

$$\frac{3}{4} N g^2 \int \frac{d^4 q}{(2\pi)^4}\frac{1}{q^4 + m^2 q^2 + \gamma^4} = 1 \ , \tag{3.27}$$

where the term $1/\eta_0$ in (3.25) has been neglected in the thermodynamic limit. The gap equation (3.27) defines the parameter γ. Notice that the dynamical mass m appears explicitly in eq.(3.27). Moreover, (3.27) reduces to the original gap equation of [18, 23] for $m = 0$. To obtain the gauge propagator, we can now go back to the expression for $\mathcal{Z}_{\rm quadr}$ which, after substituting the saddle point value $\eta = \eta_0$, becomes

$$\mathcal{Z}_{\rm quadr} = \mathcal{N} \int DA e^{-\frac{1}{2}\sum_q A^a_\mu(q)\mathcal{Q}^{ab}_{\mu\nu}A^b_\nu(-q)} \ , \tag{3.28}$$

with

$$\mathcal{Q}^{ab}_{\mu\nu} = \left(\left(q^2 + m^2 + \frac{\gamma^4}{q^2}\right)\delta_{\mu\nu} + \left(\frac{1}{\alpha} - 1\right) q_\mu q_\nu \right)\delta^{ab} \ . \tag{3.29}$$

Computing the inverse of $\mathcal{Q}^{ab}_{\mu\nu}$ and taking the limit $\alpha \to 0$, we get the gauge propagator in the presence of the dynamical gluon mass m, *i.e.*

$$\left\langle A^a_\mu(q) A^b_\nu(-q)\right\rangle = \delta^{ab}\frac{q^2}{q^4 + m^2 q^2 + \gamma^4}\left(\delta_{\mu\nu} - \frac{q_\mu q_\nu}{q^2}\right) \ . \tag{3.30}$$

Notice that, the presence of the mass m in eq.(3.30) enforces the infrared suppression of the gluon propagator.

4 The infrared behavior of the ghost propagator.

Let us discuss now the infrared behavior of the ghost propagator, given by eq.(3.15) upon contraction of the gauge fields, namely

$$\mathcal{G} \approx \frac{1}{k^2}\frac{1}{1-\rho(k)} \ , \tag{4.31}$$

with

$$\begin{aligned}\rho(k) &= \frac{g^2}{3}\frac{N}{N^2-1}\frac{k_\mu k_\nu}{k^2}\int\frac{d^4q}{(2\pi)^4}\frac{1}{(k-q)^2}\left\langle A^a_\lambda(q)A^a_\lambda(-q)\right\rangle\left(\delta_{\mu\nu}-\frac{q_\mu q_\nu}{q^2}\right) \\ &= g^2N\frac{k_\mu k_\nu}{k^2}\int\frac{d^4q}{(2\pi)^4}\frac{1}{(k-q)^2}\frac{q^2}{q^4+m^2q^2+\gamma^4}\left(\delta_{\mu\nu}-\frac{q_\mu q_\nu}{q^2}\right) \ .\end{aligned} \tag{4.32}$$

From the gap equation (3.27), it follows

$$Ng^2\int\frac{d^4q}{(2\pi)^4}\frac{1}{q^4+m^2q^2+\gamma^4}\left(\delta_{\mu\nu}-\frac{q_\mu q_\nu}{q^2}\right) = \delta_{\mu\nu} \ , \tag{4.33}$$

so that

$$1-\rho(k) = Ng^2\frac{k_\mu k_\nu}{k^2}\int\frac{d^4q}{(2\pi)^4}\frac{k^2-2qk}{(k-q)^2}\frac{1}{q^4+m^2q^2+\gamma^4}\left(\delta_{\mu\nu}-\frac{q_\mu q_\nu}{q^2}\right) \ . \tag{4.34}$$

Notice that the integral in the right hand side of eq.(4.34) is convergent and nonsingular at $k=0$. Therefore, for $k\approx 0$,

$$\left(1-\rho(k)\right)_{k\approx 0} \approx \frac{3Ng^2\mathcal{J}}{4}k^2 \ , \tag{4.35}$$

where $\mathcal{J}$ stands for the value of the integral

$$\mathcal{J} = \int\frac{d^4q}{(2\pi)^4}\frac{1}{q^2(q^4+m^2q^2+\gamma^4)} \ . \tag{4.36}$$

Finally, for the ghost propagator we get

$$\mathcal{G}_{k\approx 0} \approx \frac{4}{3Ng^2\mathcal{J}}\frac{1}{k^4} \ , \tag{4.37}$$

exhibiting the characteristic infrared enhancement which, thanks to the gap equation (3.27), turns out to hold also in the presence of the dynamical mass generation.

5 Conclusion.

In this letter we have analyzed the infrared behavior of the gluon and ghost propagators in the presence of dynamical mass generation in the Landau gauge. The restriction of the domain of integration to the Gribov region Ω has been implemented by repeating Gribov's procedure [18], which amounts to impose a no-pole condition for the two-point ghost function. The output of our analysis is summarized by equations (3.27), (3.30), (4.37). Expression (3.27) is the gap equation which defines the parameter γ. Notice now that the dynamical mass m enters

explicitly the gap equation for γ. Equation (3.30) yields the gauge propagator, which exhibits the infrared suppression. Finally, equation (4.37) establishes the enhancement of the ghost propagator. This behavior of the gluon and ghost propagators is in agreement with that found in [18, 23, 24, 25, 26, 27]. Also, lattice simulations [8, 19, 20, 21, 22] have provided confirmations of the infrared suppression of the gluon propagator and of the ghost enhancement, in the Landau gauge.

Concerning now the Gribov region Ω, it is known that it is not free from Gribov copies [30, 31, 32]. The uniqueness of the gauge condition should be ensured by restricting the domain of integration to a smaller region in field space, known as the fundamental modular region. However, this is a difficult task. Nevertheless, the restriction to the Gribov region Ω captures nontrivial nonperturbative aspects of the infrared behavior of the theory, as expressed by the suppression and the enhancement of the gluon and ghost propagators. Recently, it has been argued in [27] that the additional copies present in the Gribov region Ω might have no influence on the expectation values.

Although being outside of the aim of the present letter, we remark that the gap equation (3.27) can be also derived by using as starting point the local renormalizable action implementing the Gribov horizon, proposed in [23] by Zwanziger. It turns out in fact that the local operator $A^a_\mu A^a_\mu$ can be added to the Zwanziger action without spoiling its renormalizability [33]. This will allow to study the condensation of the operator $A^a_\mu A^a_\mu$ when the restriction to the horizon is taken into account. In this case, the combination of the algebraic BRST technique with the local composite operator formalism, see e.g. [5, 10, 29], should make possible to include the renormalization effects on the gluon and ghost propagators and to see how well these compare with the available lattice results.

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